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DESTRUCTION OF CYCLOTRON RESONANCES IN WEAKLY COLLISIONAL, INHO--ETC(U)

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NRL Memorandum Report 3916

Destruction of Cyclotron Resonances in Weakly Collisional, Inhomogeneous Plasmas

J. D. HUBA

Science Applications, Inc.
McLean, Va. 22101

and

S. L. OSSAKOW

Geophysical & Plasma Dynamics Branch
Plasma Physics Division

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April 10, 1979

This research was sponsored partially by the Defense Nuclear Agency under Subtask I25AAXYX960, Work Unit Code 17, Title "Scintillation Description"; and partially by the Office of Naval Research.

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(9) Memorandum rept.,

SECURITY CLASSIFICATION OF THIS PAGE (When Data Entered)

REPORT DOCUMENTATION PAGE

READ INSTRUCTIONS BEFORE COMPLETING FORM

1. REPORT NUMBER NRL Memorandum Report 3916	2. GOVT ACCESSION NO.	3. RECIPIENT'S CATALOG NUMBER

4. TITLE (and Subtitle) (6) DESTRUCTION OF CYCLOTRON RESONANCES IN WEAKLY COLLISIONAL, INHOMOGENEOUS PLASMAS .	5. TYPE OF REPORT & PERIOD COVERED Interim report on a continuing NRL problem 6. PERFORMING ORG. REPORT NUMBER
7. AUTHOR(s) (10) J. D. Huba [redacted] S. L. Ossakow	8. CONTRACT OR GRANT NUMBER(s)
9. PERFORMING ORGANIZATION NAME AND ADDRESS Naval Research Laboratory Washington, D.C. 20375	10. PROGRAM ELEMENT, PROJECT, TASK AREA & WORK UNIT NUMBERS NRL Problems 67H02-42D and 67A03-16B DNA Subtask I25AAXYX960
11. CONTROLLING OFFICE NAME AND ADDRESS Defense Nuclear Agency, Washington, D.C. 20305 & Office of Naval Research, Arlington, VA 22217. (11)	12. REPORT DATE 1Ø Apr[redacted]79 13. NUMBER OF PAGES 27
14. MONITORING AGENCY NAME & ADDRESS (if different from Controlling Office) (12) 28p.	15. SECURITY CLASS. (of this report) UNCLASSIFIED 15a. DECLASSIFICATION/DOWNGRADING SCHEDULE N/A

16. DISTRIBUTION STATEMENT (of this Report)

Approved for public release; distribution unlimited.

(14) NRL-MR-3916

17. DISTRIBUTION STATEMENT (of the abstract entered in Block 20, if different from Report)

(16) I25AAXY (17) X96Ø

18. SUPPLEMENTARY NOTES

This research was sponsored partially by the Defense Nuclear Agency under Subtask I25AAXYX960, Work Unit Code 17, Title "Scintillation Description"; and partially by the Office of Naval Research.

19. KEY WORDS (Continue on reverse side if necessary and identify by block number)

Collisional Kinetic Drift Waves
Cyclotron Resonances
Ion Cyclotron Drift Instability
Lower Hybrid Drift Instability
Equatorial Spread F Application
Tandem Mirror Experiment Application

omega sub rho   Omega

20. ABSTRACT (Continue on reverse side if necessary and identify by block number)

It is shown, both analytically and numerically, that cyclotron resonances can be destroyed in dense ($\omega_p > \Omega$, where ω_p is the plasma frequency and Ω is the cyclotron frequency), weakly collisional, inhomogeneous plasmas when $(\nu/\Omega)k^2 r_L^2 \gtrsim 1$, where ν is the collision frequency and r_L is the mean Larmor radius. The theory is based upon a model Fokker-Planck equation. It is found that the particles make a transition from magnetized to unmagnetized behavior. This is an important result since it indicates that the ion- and electron-cyclotron-drift instabilities transform into

nu   sub L

(nu/Omega) k-squared (r sub L) squared somewhat > (Continued) next page

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20. ABSTRACT (Continued)

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their unmagnetized counterparts, the lower-hybrid-drift instability and the ion acoustic instability, respectively. The ion-cyclotron-drift instability (or drift-cyclotron instability) is examined in detail and is found to become the lower-hybrid-drift-instability in the region of maximum growth when $(m_e/m_i)^{1/2} \lesssim \omega/\Omega_i \gtrsim \nu_{ii}/\Omega_i \gtrsim m_e/m_i$ for $T_e \approx T_i$ plasmas. The first inequality is required to overcome electron viscous damping, while the second allows the ions to become unmagnetized. Applications to the equatorial F region of the ionosphere and the Tandem Mirror Experiment (TMX) are discussed.

sub i

sub e

approximately equal

somewhat >

Omega/Omega sub i somewhat > nu sub ii / Omega sub i

- sq.rt.(m sub e/m sub i)

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DESTRUCTION OF CYCLOTRON RESONANCES IN WEAKLY COLLISIONAL, INHOMOGENEOUS PLASMAS

I. INTRODUCTION

We investigate the effect of like particle collisions (i.e., ion-ion and electron-electron collision) on short wavelength, electrostatic cyclotron instabilities in weakly collisional, inhomogeneous plasmas. That is, for $\nu/\Omega << 1$ and $k_\perp^2 r_L^2 >> 1$ where ν is the collision frequency, Ω is the cyclotron frequency and r_L is the mean Larmor radius. We base our analysis on a model Fokker-Planck equation originally proposed by Chandresekhar[1] and applied to a homogeneous, unmagnetized plasma by Lenard and Bernstein[2] and to a homogeneous, magnetized plasma by Dougherty.[3] Dougherty found that cyclotron resonances could be destroyed by collisions if $\nu/\Omega \gtrsim 1$ or $(\nu/\Omega) k_\perp^2 r_L^2 \gtrsim 1$. The latter condition is particularly interesting since it indicates that cyclotron waves with $k_\perp^2 r_L^2 >> 1$ may not exist even though $\nu/\Omega << 1$. Physically this occurs because the particles can diffuse a distance $L_D \sim (\nu/\Omega)^{1/2} r_L$ in one gyroperiod. If this distance is greater than a wavelength (i.e., $(\nu/\Omega) k_\perp^2 r_L^2 \gtrsim 1$) then the cyclotron wave cannot maintain its coherence. In fact, it has been shown analytically that the electron-Bernstein-mode dispersion equation makes a transition to the ion-acoustic-mode dispersion equation (i.e., the electrons become "unmagnetized") when $(\nu_{ee}/\Omega_e) k_\perp^2 r_{Le}^2 \geqslant 1$.[4] We mention that this case has been studied numerically and such a transition was not observed.[5]

We demonstrate in this paper that in weakly collisional, inhomogeneous plasmas, the ion- and electron-cyclotron drift instabilities transform into their unmagnetized counterparts, the lower-hybrid-drift and the ion acoustic instabilities, respectively. We discuss the ion-cyclotron-drift instability (also known as the drift-cyclotron instability)[6] in detail and show it becomes the

Manuscript Submitted January 31, 1979

lower-hybrid-drift instability[7] in the region of maximum growth when $(m_e/m_i)^{1/2}$ $(\omega/\Omega_i) \gtrsim$ $\nu_{ii}/\Omega_i \gtrsim m_e/m_i$ for $T_e \approx T_i$ plasmas. The first inequality is required to overcome electron viscous damping, while the second allows the ions to become "unmagnetized." Applications to the equatorial Spread F phenomenon[8] and the Tandem Mirror Experiment (TMX)[9] are discussed.

We point out that several Russian authors have also studied the effects of collisions on instabilities in inhomogeneous plasmas. Mikhailovskii and Pogutse[10] have derived a general dispersion tensor, including electrostatic and electromagnetic perturbations, based on the Bhatnagar-Gross-Krook (BGK) collision model.[11] Rukhadze and Silin[12] base their analysis on the Landau collision integral and discuss a variety of drift instabilities. However, neither of the above studies treat the limit $(\nu/\Omega)k_\perp^2 r_L^2 \gtrsim 1$ and hence, do not find a transition to "unmagnetized" behavior.

The structure of the paper is as follows. In the next section we derive a general dispersion equation for electrostatic waves in a weakly collisional, inhomogeneous plasma based upon a model Fokker-Planck equation. In Section III we apply this theory to the ion-cyclotron-drift instability, presenting both analytical and numerical results. Finally, in the last section we discuss the implications of these results for both space and laboratory plasmas.

II. THEORY

A. Physical Configuration and Assumptions

The physical configuration which we consider is described as follows. The equilibrium magnetic field is $\mathbf{B} = B_0\hat{e}_z$, the density varies only in the x direction and the temperature is assumed constant. The density gradient produces a cross-field drift given by $\mathbf{U}_0 = V_d\,\hat{e}_y$ where $V_d = (v_{th}^2/2\Omega)\partial \ln n/\partial x$ is the diamagnetic drift velocity. Here, $v_{th}^2 = 2T/m$ is the thermal velocity and $\Omega = qB_0/mc$ is the cylcotron frequency for a species with temperature T, charge q and mass m. We treat only electrostatic oscillations and assume perturbation quantities to vary as $\exp[-i(\mathbf{k}\cdot\mathbf{x} - \omega t)]$ where $\mathbf{k} = k_\perp \hat{e}_y + k_\parallel \hat{e}_z$. We make use of the local approximation which requires $kL_n >> 1$ where $L_n = (\partial \ln n/\partial x)^{-1}$ is the scale length of the density inhomogeneity. The equilibrium distribution function used in the analysis is[13]

$$f_0 = n_0(x)\left(\frac{1}{\pi v_{th}^2}\right)^{3/2} \exp\left[-(\mathbf{v} - \mathbf{U}_0)^2/v_{th}^2\right]. \tag{1}$$

where we emphasize that this is a self-consistent Vlasov equilibrium for strongly inhomogeneous plasmas (i.e., $L_n \gtrsim r_L$ where r_L is the mean gyroradius) locally, say at $x = x_0$.

B. Model Fokker-Planck Equation

We choose the following model Fokker-Planck equation to describe collisional effects[1-3]

$$\frac{\partial f}{\partial t} + v_j\frac{\partial f}{\partial x_j} + \frac{q}{m}\left(\mathbf{E} + \frac{1}{c}\mathbf{v}\times\mathbf{B}\right)_j\frac{\partial f}{\partial v_j} = \nu\frac{\partial}{\partial v_j}\left[(v_j - U_j)f + \frac{T}{m}\frac{\partial f}{\partial v_j}\right] \tag{2}$$

where $n = \int d^3v f$, $nV_i = \int d^3v\, v_i f$ and $3nT = \int d^3v\, m(\mathbf{v}-\mathbf{V})^2 f$. The first term in brackets represents friction and is taken to be proportional to the velocity relative to the mean velocity. This is a reasonable approximation for thermal particles but is incorrect for "tail" particles. The second term represents diffusion and is chosen to be isotropic which requires the Debye length to be less than the mean Larmor radius. This model describes small-angle collisions and is particularly useful in handling like particle collisions.

Linearizing Eq. (2) we obtain

$$\frac{\partial f_1}{\partial t} + v_j\frac{\partial f_1}{\partial x_j} + \left[\frac{q}{m}\left(\frac{1}{c}\mathbf{v}\times\mathbf{B}_o\right)_j - \nu(v_j - U_{jo})\right]\frac{\partial f_1}{\partial v_j} - 3\nu f_1 - \nu\frac{v_{th}^2}{2}\frac{\partial^2 f_1}{\partial v_j^2} = h(\mathbf{x},\mathbf{v},t) \tag{3}$$

where

$$h(\mathbf{x},\mathbf{v},t) = \frac{q}{m}\frac{\partial\phi_1}{\partial x_j}\frac{\partial f_o}{\partial v_j} + \nu\frac{\partial}{\partial v_j}\left[-u_j f_o + \frac{T_1}{m}\frac{\partial f_o}{\partial v_j}\right]. \tag{4}$$

$n_1 = \int d^3v f_1, n_o u_j = \int d^3v\, v_j f_1, 3n_o(T_1/T_o) = (2/v_{th}^2)\int d^3v(\mathbf{v} - \mathbf{U}_o)^2 f_1$ and $v_{th}^2 = 2To/m$. Here, the subscripts o and 1 denote unperturbed and perturbed quantities, respectively, and $U_j = U_{jo} + u_j$.

We can cast Eq. (3) into a form which can be solved using a Green's function[3] by the following transformation: $\mathbf{x} \to \mathbf{x} + \mathbf{V}t$ and $\mathbf{v} \to \mathbf{v} + \mathbf{V}$ where

$$V_x = \frac{\Omega\nu}{\nu^2 + \Omega^2} U_o; \quad V_y = \frac{\nu^2}{\nu^2 + \Omega^2} U_o.$$

Making use of this coordinate change we find that Eq. (3) becomes

$$\frac{\partial f_1}{\partial t} + v_j \frac{\partial f_1}{\partial x_j} + \left[\frac{q}{m}\left(\frac{1}{c}\,\mathbf{v} \times \mathbf{B}_o\right)_j - \nu v_j\right]\frac{\partial f_1}{\partial v_j} - 3\nu f_1 - \nu \frac{v_{th}^2}{2}\frac{\partial^2 f_1}{\partial v_j^2} = h(\mathbf{x}, \mathbf{v}, t) \qquad (5)$$

and h is defined by Eq. (4) with the transformation reflected in f_o. We point out that the above coordinate transformation is sign dependent (through the diamagnetic drift velocity U_o) and *cannot be made simultaneously* for both species. However, this does ot create a serious problem since we only consider collisional effects on a single species. Moreover, since we are interested in the situation where $\nu/\Omega << 1$ and $k_\perp^2\, r_L^2 >> 1$ we note that (1) the old reference frame is a good approximation to the new one and (2) the final term in Eq. (4) can be neglected.[4] Thus, within the context of these assumptions, we consider

$$\frac{\partial f_1}{\partial t} + v_j \frac{\partial f_i}{\partial x_j} + \left[\frac{q}{m}\left(\frac{1}{c}\,\mathbf{v} \times \mathbf{B}_o\right)_j - \nu v_j\right]\frac{\partial f_1}{\partial v_j} - 3\nu f_1 - \nu \frac{v_{th}^2}{2}\frac{\partial^2 f_1}{\partial v_j^2} = \frac{q}{m}\frac{\partial \phi_1}{\partial x_j}\frac{\partial f_o}{\partial v_j} \qquad (6)$$

to describe electrostatic oscillations in a weakly collisional, inhomogeneous plasma where f_o is defined by Eq. (1). The important physics contained in Eq. (6) is the modification of the unperturbed orbits to include a resistive term and the diffusion of perturbations as they are connected along an unperturbed orbit.

C. Derivation of the Dispersion Equation

The dispersion equation for electrostatic perturbations is

$$D(\omega, k) = 1 + \sum_\alpha \chi_\alpha = 0 \qquad (7)$$

where χ_α is the susceptability of species α is defined by $\chi_\alpha \phi_1 = -(4\pi e_\alpha/k^2) n_{\alpha 1}$. Thus, we only need to know the perturbed density in order to determine the linear behavior of electrostatic waves. In order to compute the perturbed density we follow the method of Dougherty and employ a Green's function.[3] We only present the results here and refer the interested reader to Ref. 3 for the mathematical details. In the absence of collisions, the Green's function is a δ function and perturbations are simply convected along the unperturbed orbit. When collisions are included, the perturbations also spread out due to the diffusion as well as being convected along the unperturbed orbit (which now includes a resistive term). Thus, the Green's function becomes a Gaussian distribution in phase space.

Using notation similar to Dougherty, we find

$$n_1 = \frac{2}{v_{th}^2}\, a^\mu\, H_\mu \tag{8}$$

where

$$H_\mu = \left[i \frac{\partial}{\partial \rho^\mu} I \right]_{\sigma=\rho=o}, \tag{9}$$

$$I = \frac{n_o}{\Omega} \int_o^\infty d\tau \exp[i\tilde{\omega}\tau - \Phi(\tau) - \Psi(\tau)], \tag{10}$$

$$\begin{aligned}\Phi(\tau) = {} & \frac{1}{2}\, k_\perp^2\, r_L^2\, [\cos\theta + \tilde{\nu}\tau - \exp(-\tilde{\nu}\tau)\cos(\tau - \theta)] \\ & + \frac{1}{2} \frac{k_\parallel^2\, r_L^2}{\tilde{\nu}^2} [\tilde{\nu}\tau - 1 - \exp(-\tilde{\nu}\tau)] \\ & + i\frac{k_\perp U_o}{\Omega} [\exp(-\tilde{\nu}\tau)\sin\tau + \tilde{\nu}(1 - \exp(-\tilde{\nu}\tau)\cos\tau)],\end{aligned} \tag{11}$$

$$\begin{aligned}\Psi(\tau) = \frac{v_{th}^2}{2} \Bigg\{ & \frac{1}{2}\sigma_\lambda\, \sigma^\lambda + \frac{1}{2}\rho_\lambda\, \rho^\lambda + \frac{k_\lambda\, \sigma^\lambda}{\Omega}\, \frac{1 - \exp[-(\tilde{\nu} - i\lambda)\tau]}{\tilde{\nu} - i\lambda} \\ & + \frac{k_\lambda\, \rho^\lambda}{\Omega}\, \frac{1 - \exp[-(\tilde{\nu} + i\lambda)\tau]}{\tilde{\nu} + i\lambda} + \sigma_\lambda\, \rho^\lambda \exp[-(\tilde{\nu} + i\lambda)\tau] \\ & + \frac{2i\sigma_\lambda\, U_o^\lambda}{U_{th}^2} \exp[-(\tilde{\nu} + i\lambda)\tau] \Bigg\},\end{aligned} \tag{12}$$

$$a^\mu = -\frac{q}{m}\phi_1 (ik^\mu - \epsilon_n^\mu),$$

$\theta = 2\tan^{-1}\tilde{\nu}, \tilde{\nu} = \nu/\Omega, \tilde{\omega} = \omega/\Omega$ and $\epsilon_n^\mu = (\partial \ln n/\partial x_\mu)_{x=x_o}$. Note that we use a polarized coordinate system (x^1, x^0, x^{-1}) where $x^1 = (1/\sqrt{2})(x + iy), x^0 = z$ and $x^{-1} = (1/\sqrt{1})(x - iy)$. This transformation is unitary but not orthogonal so we must distinquish contravariant and covariant vectors.

Making use of Eqs. (8) - (12), we arrive at

$$\chi = \frac{2\omega_p^2}{k^2 v_{th}^2}\left\{1 + i\left\{\frac{\omega - k_\perp U_o}{\Omega}\right\}\int_o^\infty d\tau \exp[i\tilde{\omega}\tau - \Phi(\tau)]\right\} \tag{13}$$

where $\omega_p^2 = 4\pi nq^2/m$. We emphasize that this expression for χ is only valid for $\nu/\Omega << 1$ and $k_\perp^2 r_L^2 >> 1$. Note that in the limit $L_n \rightarrow \infty$ (i.e., $U_o \rightarrow 0$), Eq. (13) reduces to a result equivalent to that derived by Allan and Sanderson[4] and in the collisionless limit $(\tilde{\nu} = o)$ becomes

$$x = \frac{2\omega_p^2}{k^2 v_{th}^2}\left[1 + i\left[\frac{\omega - k_\perp U_o}{\Omega}\right]\int d\tau \exp\left\{i\omega\tau + \frac{1}{2}k_\perp^2 r_L^2(\cos\tau - 1) - i\frac{k_\perp U_o}{\Omega}\sin\tau + \frac{1}{4}k_\parallel r_L^2\tau^2\right\}\right] \tag{14}$$

which is a standard result.[15]

III. APPLICATION: ION-CYCLOTRON-DRIFT INSTABILITY

The ion-cyclotron-drift (ICD) instability (also known as the drift-cyclotron instability)[6] has received considerable attention in the past 15 years. It is thought to play an important role as an anomalous transport mechanism in a variety of "collisionless" plasmas (e.g., post-implosion stage of a theta pinch,[16] solenoidal section of TMX,[17] the polar cusp,[18] the ionosphere,[19] the Earth's magnetotail[20]). The instability can be excited when $L_n/r_{Li} < (m_i/m_e)^{1/2}/l$ and is characterized by $\omega_r \sim l\Omega_i$ and $\gamma \sim l\Omega_i(m_e/m_i)^{1/4}$ where $\omega = \omega_r + i\gamma$ and l is the ion cyclotron harmonic number. Maximum growth occurs for $\mathbf{k} \cdot \mathbf{B} = 0$ (i.e., $k_\parallel = 0$) with

$kr_{Le} \sim (T_e/T_i)^{1/2}$ in dense plasmas (i.e., $\omega_{pe}^2 >> \Omega_e^2$). Note that to satisfy the condition $(\nu_{ii}/\Omega_i)k_\perp^2 r_{Li}^2 \gtrsim 1$ at maximum growth, we require

$$\frac{\nu_{ii}}{\Omega_i} \gtrsim \frac{m_e}{m_i}. \tag{15}$$

For a hydrogen plasma, Eq. (15) implies that if $(\nu_{ii}/\Omega_i) \gtrsim 5 \times 10^{-4}$, ion-ion collisions will destroy the cyclotron resonances and hence, the ICD instability. However, rather than drive the plasma to stability we demonstrate, both analytically and numerically, the transformation of the ICD instability to the lower-hybrid-drift (LHD) instability in the presence of weak ion-ion collisions.

The dispersion equation under consideration is

$$D(\omega,\ k) = 1 + \chi_i + \chi_e = 0 \tag{16}$$

where

$$\chi_i = \frac{2\omega_{pi}^2}{k^2 v_i^2}\left[1 + i\left(\frac{\omega - k_\perp V_{di}}{\Omega_i}\right)\int_o^\infty d\tau \exp(i\omega\tau/\Omega_i - \Phi_i(\tau))\right], \tag{17}$$

$$\chi_e = \frac{2\omega_{pe}^2}{k^2 v_e^2}\left[1 - \exp(b_e) I_o(b_e)\left(1 - \frac{k_\perp V_{de}}{\omega}\right)\right], \tag{18}$$

and

$$\Phi_i(\tau) = \frac{1}{2} k_\perp^2\ r_{Li}^2\left[\cos\sigma + \tilde{\nu}_i\tau - \exp(-\tilde{\nu}_i\tau)\cos(\tau-\theta)\right] \tag{19}$$
$$+ i\frac{k_\perp V_{di}}{\Omega_i}\left[\exp(-\tilde{\nu}_i\tau)\sin\tau + \tilde{\nu}_i(1 - \exp(-\tilde{\nu}_i\tau)\cos\tau)\right],$$

$\tilde{\nu}_i = \nu_{ii}/\Omega_i$, $b_e = \frac{1}{2}\ k^2\ r_{Le}^2$, $\omega_{p\alpha}^2 = 4\pi n e_\alpha^2/m_\alpha$, $v_\alpha^2 = 2T_\alpha/m_\alpha$, $V_{d\alpha} = (v_\alpha^2/2\Omega_\alpha)(\partial \ln n/\partial x)_{x=x_0}$

$\Omega_\alpha = e_\alpha B_0/m_\alpha c$, I_l is the modified Bessel function of order l. The simple form chosen for the electron susceptibility is valid since $\Omega_i \lesssim \omega << \Omega_e$ and it is assumed that $\nu_{ei}/\Omega_e << 1$, $\nu_{ee}/\Omega_e << 1$ and $\nu_{ee}k^2\ r_{Le}^2 << \omega$ (i.e., the electrons are collisionless). The final condition is required to avoid collisional damping due to electron viscosity[21] and will be discussed further in the final section. We emphasize that Eq. (17) is valid only in the regime $\nu_{\shortparallel}/\Omega_i << 1$ and $k_\perp^2 r_{Li}^2 >> 1$.

A. Analytical Results

i. $\tilde{\nu}(k^2 r_{Li}^2) << 1$

In the limit $\tilde{\nu}(k^2 r_{Li}^2) << 1$, $\tilde{\nu} << 1$ and $k^2 r_{Li}^2 >> 1$ collisional effects can be ignored and we find that

$$\chi_i = \frac{2\omega_{pi}^2}{k^2 v_i} \left[1 - \frac{\omega - kV_{di}}{\Omega} \frac{\sqrt{\pi}}{kr_{Li}} \cot(\pi\tilde{\omega})\right] \tag{20}$$

and

$$\chi_e = \frac{2\omega_{pe}^2}{k^2 v_e^2} \left[1 - \left(1 - \frac{kV_{de}}{\omega}\right) \exp(-b_e) I_0(b_e)\right] \tag{21}$$

which are the standard results for the ion-cyclotron-drift instability. Noting that $\pi \cot(\pi z) = \sum_{m=-\infty}^{\infty} \frac{1}{z-m}$, we can rewrite Eq. (20) in a more conventional form[6]

$$\chi_i = \frac{2\omega_{pi}^2}{k^2 v_i^2} \left[1 - \frac{1}{\sqrt{\pi}} \frac{1}{kr_{Li}} \sum_m \frac{\omega - kV_{di}}{\omega - m\Omega}\right] \tag{22}$$

where $(\sqrt{\pi} k r_{Li})^{-1}$ is the large argument expansion of $I_m \left(\frac{1}{2} k^2 r_{Li}^2\right) \exp\left(-\frac{1}{2} k^2 r_{Li}^2\right)$. The ion cyclotron resonances are apparent in Eq. (22).

ii. $\tilde{\nu}(k^2 r_{Li}^2) \gtrsim 1$

In the limit $\tilde{\nu}(k^2 r_{Li}^2) \gtrsim 1$, $\tilde{\nu} << 1$ and $k^2 r_{Li}^2 >> 1$ we can again approximate the electron susceptibility by Eq. (21). However, the ion susceptibility is changed substantially from Eq. (20). We consider Eqs. (17) and (19) and point out that for $k_\perp^2 r_{Li}^2 >> 1$ the dominant contribution to the integrand occurs in the vicinity of $\tau \approx 2\pi m (m = 0, 1, 2, ..)$. We let $\tau = 2\pi m + t$ and obtain

$$\int_0^\infty d\tau \exp(i\tilde{\omega}\tau - \Phi_i(\tau)) \approx \int_0^{\delta_0} dt \exp\left[i\left(\frac{\omega - kV_{di}}{\Omega}\right)t - \frac{1}{2}k^2r_{L_i}^2t^2\right] \tag{23}$$

$$+ \sum_{m=1}^{\infty} \int_{2\pi m-\delta_m}^{2\pi m+\delta_m} dt \exp\left[i\left(\frac{\omega - kV_{di}}{\Omega}\right)t - \frac{1}{2}k^2r_{L_i}^2t^2\right] \exp[-2\pi k^2 r_{L_i}^2 m\tilde{\nu}] \exp\left[+2\pi im\left(\frac{\omega - kV_{di}}{\Omega_i}\right)\right]$$

where δ_0, δ_m << 1. Clearly when $\tilde{\nu}(k^2r_{L_i}^2) \gtrsim 1$, the first term in Eq. (23) is dominant and we find (letting $\delta_0 \to \infty$) that

$$\chi_i \approx \frac{2\omega_{p_i}^2}{k^2 v_i^2}[1 + \zeta_i Z(\zeta_i)] \tag{24}$$

where $\zeta_i = (\omega - kV_{di})/kv_i$ and we have made use of the relation

$$Z(\zeta) = i\sqrt{\pi}\exp(-\zeta^2)[1 + erf(i\zeta)]. \tag{25}$$

Thus, the dispersion equation becomes

$$D(\omega, k) = 1 + \frac{2\omega_{pi}^2}{k^2 v_i}[1 + \zeta_i Z(\zeta_i)] + \frac{2\omega_{pe}^2}{k^2 v_e^2}\left[1 - \left(1 - \frac{kV_{de}}{\omega}\right)\exp(-b_e)I_0(b_e)\right] \tag{26}$$

which describes the lower-hybrid drift instability.[7] Physically, the instability is excited since the ions can now move across the field lines and be in resonance with the drift wave. We mention that this analysis parallels that of Allan and Sanderson who considered electron Bernstein modes in a homogeneous plasma.[22] Finally, note from Eq. (23) that the ICD instability will also transform into the LHD instability in the strong drift velocity regime (i.e. $\gamma > \Omega_i$).

B. Numerical Analysis

We now solve Eqs. (15)-(18) numerically. As a relevant example, we consider the equatorial F region of the ionosphere where the dominant ionic component is O^+ so that $m_i = 16\ m_p$. Recent experimental observations indicate intense VHF and UHF radar backscatter during equatorial spread F resulting from density irregularities of 1 m and 36 cm, respectively.[19] These irregularities correspond to wavelengths such that $k^2r_{L_i}^2 >> 1$. Since there is also evidence of density inhomogeneities with scale lengths $L_n/r_{Li} \lesssim (m_i/m_e)^{1/2}$,[23] it has been suggested that the ICD or LHD instabilities are responsible for the irregularities. Also, because $\nu_{ii}/\Omega_i < 10^{-2}$ ion-ion collisions can play a significant role and a numerical study is required.

In Fig. 1 we show γ/Ω_i vs. kr_{Le} and kr_{Li} for typical ionospheric parameters: $T_e = T_i$, $\omega_{pe}/\Omega_e = 10.0$, $V_{di}/v_i = 0.037$ which corresponds to $L_n \approx 13\, r_{Li} (\approx 75$ m) and $\nu_{ii}/\Omega_i = 0.0$, 10^{-6}, 10^{-5}, 10^{-4}. We comment that wave growth also occurs for $kr_{Le} > 1.5$ but has not been plotted. We first indicate the wave growth spectrum for the LHD instability (—) where we have used Eq. (25). The spectrum is broadband with maximum growth at $kr_{Le} \approx 1.2$ with $\omega_r \approx 3.0\Omega_i$ and $\gamma \approx 0.08\ \Omega_i$. Growth occurs for $kr_{Le} < 0.15$ but is very weak (i.e., $\gamma < 10^{-3}\Omega_i$). For $\nu_{ii}/\Omega_i = 0.0$ (— · — · —) we see that the first three cyclotron harmonics of the ICD instability are excited with maximum growth occurring at $kr_{Le} \approx 1.0$ with $\omega_r \approx 3.0\ \Omega_i$ and $\gamma \approx 0.15\ \Omega_i$. As ν_{ii}/Ω_i is increased to 10^{-6} (— — — —) we already note a change in the wave spectrum. Although the harmonic structure is maintained, the maximum growth rate decreases 6% and growth is extended beyond $kr_{Le} \sim 1.45$ for the 3rd harmonic. For $\nu_{ii}/\Omega_i = 10^{-5}$ (······) a more dramatic change occurs in the wave spectrum as the harmonic structure becomes distorted. The wave spectrum collapses to the LHD curve for $kr_{Le} \approx 1.4$ ($kr_{Li} \approx 210$) which corresponds to $(\nu_{ii}/\Omega_i)k^2 r_{Li}^2 \sim 0.5$. Moreover the maximum growth rate decreases by 40% to $\gamma \sim 0.09\ \Omega_i$. The first two harmonics are still evident although their growth rates have decreased somewhat. When $\nu_{ii}/\Omega_i = 10^{-4}$ (xxxxx) the wave spectrum falls on the LHD curve and only the slighest hint of harmonic structure remains at $l = 1$ and $l = 2$.

Figure 2 plots γ_M/Ω_i vs. V_{di}/v_i for $T_e = T_i$, $\omega_{pe}/\Omega_e = 10.0$ and $\nu_{ii}/\Omega_i = 0$, 10^{-6}, 10^{-5}, 10^{-4}. Here γ_M is the growth rate maximized with respect to k. Again we indicate the results of the LHD instability (—) as a reference. Setting $\nu_{ii}/\Omega_i = 0$ we recover the ICD instability which displays substantially stronger growth than the LHD instability for weak drifts (i.e., $V_{di} < 0.10\ v_i$). We can distinguish the excitation of the first 8 harmonics as V_{di}/v_i is increased. Also, note the transition to the LHD instability for $V_{di} > 0.11\ \Omega_i$ which occurs when $\gamma \sim 0.75\ \Omega_i$ as indicated earlier. For $\nu_{ii}/\Omega_i = 10^{-6}$ (— — — —) we note that the growth rates are generally reduced for $V_{di} < 0.085\ v_i$ although 8 harmonics are still apparent. Increasing ν_{ii}/Ω_i to 10^{-5} (·····) causes the wave spectrum to collapse to the LHD spectrum with only the first 5 harmonics visible. Finally, for $\nu_{ii}/\Omega_i = 10^{-4}$ the curve lies on the LHD plot.

Thus, we have shown numerically that the ICD instability makes a transition to the LHD instability for $\frac{\nu_{ii}}{\Omega_i} k^2 r_{Li}^2 \geq 1$. Since growth occurs for $k^2 r_{Li}^2 >> 1$, we find that only a very small amount of collisions $\left(\frac{\nu_{ii}}{\Omega_i} \geq 10^{-5}\right)$ is needed to cause this transition in the region of maximum growth. For typical ionospheric conditions, $10^3\ \mathrm{cm}^{-3} \leq n \leq 10^6\ \mathrm{cm}^{-3}$, $T_i \sim 0.1$ eV and $B \sim 0.3\ G$ we find that $10^{-4} \leq \frac{\nu_{ii}}{\Omega_i} \leq 10^{-1}$ so that the LHD instability should dominate when instability occurs. However, it is found that the condition $\nu_{ee}\ k^2\ r_{Le}^2 < \omega$ is not satisfied when $n > 10^4\ \mathrm{cm}^{-3}$ for moderate density gradients ($L_n > 8.0\ r_{Li}$) and it is expected that the LHD instability will be stabilized by the viscous damping of the electrons for this situation. Hence, instability is likely to occur infrequently and will be restricted to low density regions of the F region, such as the interior of ionspheric bubbles during equatorial Spread F[24]. We defer a more complete discussion of this instability and its application to Spread F to a subsequent report.

IV. DISCUSSION

We have demonstrated, both analytically and numerically, that cyclotron resonances can be destroyed in dense ($\omega_p > \Omega$), weakly collisional, inhomogeneous plasmas when $(\nu/\Omega)k^2 r_L^2 \geq 1$. The physical mechanism for this phenomenon is that particles can diffuse a distance $L_D \sim (\nu/\Omega)^{1/2} r_L$ in one gyroperiod due to collisions. If this distance exceeds the wavelength (i.e., $(\nu/\Omega)k^2 r_L^2 \geq 1$) then the wave can no longer maintain its coherence. This result is in accordance with the work of Dougherty[3] and Allan and Sanderson[4] who considered only homogeneous plasmas. It is an important result since it indicates that the ion and electron drift-cyclotron instabilities transform into their unmagnetized counterparts, the lower-hybrid-drift (LHD) instability and the ion acoustic (IA) instability, respectively, for weakly collisional plasmas. We specifically discussed the ICD instability in detail and found the transition to occur in the region of maximum growth for $(\nu_{ii}/\Omega_i) \gtrsim (m_e/m_i)$. However, we have also pointed out that electron viscous damping can stabilize the LHD instability when $\nu_{ee}\, k^2\, r_{Le}^2 > \omega$ since electron viscosity provides a dissipative mechanism for the wave. Recognizing that $\nu_{ee} \sim (m_i/m_e)^{1/2}\, \nu_{ii}^{25}$, we find that $(m_e/m_i)^{1/2}\, \omega/\Omega_i > \nu_{ii}/\Omega_i$ is required for instability to occur in the region of maximum growth. Electron-electron collisions then have the effect of placing a threshold condition on the diamagnetic drift velocity necessary to excite the LHD instability since $\omega \sim kV_{di}$. Hence, "collisionless" plasmas on cyclotron time scales, such as the F region in the ionosphere, may in fact be strongly affected by weak collisions as we have shown.

In collisionless plasmas both the ion and electron cyclotron drift instabilities require a critical density gradient scale length to be excited. This occurs because the instabilities are produced by a coupling of a cyclotron wave $\omega_1 = \Omega$ and a drift wave $\omega_2 = kV_d$, and a critical drift velocity is therefore required to satisfy the matching condition $\omega_i \approx \omega_2$. However, when $(\nu/\Omega)k^2 r_L^2 \gtrsim 1$ these instabilities transform into the LHD and IA instabilities which are driven via wave-particle resonances and have different turn-on criteria. In the case of the LHD instability a critical drift velocity is required to overcome electron viscuous damping as discussed above. On the other hand, the critical drift velocity for the IA instability is sensitive to T_e/T_i and may or may not be excited depending on the actual plasma conditions.

The difference in the nonlinear behavior between the ICD and LHD instabilities, and the ECD and IA instabilities is probably a more important consequence of the theory since entirely different stabilization mechanisms are operable. The ECD instability has been studied in detail and it has been found that in the collisionless limit, resonance broadening due to turbulence effectively "unmagnetizes" the electrons and a transition is made to the IA instability.[26] This is the same result we have found for $(\nu_{ee}/\Omega_e)k^2 r_{Le}^2 \gtrsim 1$. A variety of theories have been proposed for the nonlinear saturation of the IA instability and they will not be discussed here.

Recent nonlinear theories have been presented for the stabilization of the ICD instability and rely upon a nonlinear frequency shift[27] or low-frequency ($\omega << \Omega_i$) density fluctuations.[28] On the other hand, the suggested saturation mechanisms for the LHD instability are quasilinear stabilization,[29] ion trapping[30] and electron resonance broadening.[31] We comment that Hasegawa's theory[28] concerning stabilization via low-frequency density fluctuations is also applicable to the LHD instability. We will not discuss the details of these stabilization mechanisms but simply emphasize that different saturation energies are anticipated for the ICD and LHD instabilities. Hence, the role of weak ion-ion collisions can substantially alter the nonlinear plasma response due to different anomalous transport properties.

Although we have only discussed an application of this theory to the ionosphere, we wish to point out its relevance to the Tandem Mirror Experiment (TMX).[9]. TMX will consist of a straight solenoidal plasma which is confined by conventional mirror end cells. The solenoidal section will be similar to a theta pinch and the ions are expected to have a roughly Maxwellian distribution. Since density gradients are expected with scale lengths such that $L_n/r_{Li} < (m_i/m_e)^{1/2}$, it is anticipated that the ICD instability is likely to occur. Using expected operating parameters for the solenoidal section ($T_i = 80$ eV, $n \approx 10^{13}$ cm^{-3}, $B = 500$ G), we find that $\nu_{ii}/\Omega_i \sim 4 \times 10^{-3}$. Thus, we conclude that drift waves with $kr_{Li} \gtrsim 15$ will probably be excited by the LHD instability and the nonlinear dynamics in this regime can be studied using unmagnetized ions. Of course, finite β and temperature gradient effects are also expected to play a role in TMX[17,32,33] which are not considered in this paper. However, in principle these effects can be included in a straightforward manner and are not expected to alter this conclusion.

Finally, we point out that the Green's function method used in this work is not the most direct technique, although it lends itself to a simple physical interpretation. Recently, Catto has computed the perturbed density in a collisional, inhomogeneous plasma using a velocity transform method[34] which is a simpler analysis. For a more complex situation (e.q., inclusion of temperature gradients, ∇B and electromagnetic effects) this procedure would be preferable to the one outlined in this work.

ACKNOWLEDGMENTS

We wish to express our sincere thanks to Dr. J.J. Sanderson for a very beneficial communication. We also thank Dr. P.K. Chaturvedi and Dr. J. Drake for several helpful discussions, and Dr. C. S. Liu for calling our attention to the importance of electron viscosity. This research has been supported by the Defense Nuclear Agency and the Office of Naval Research.

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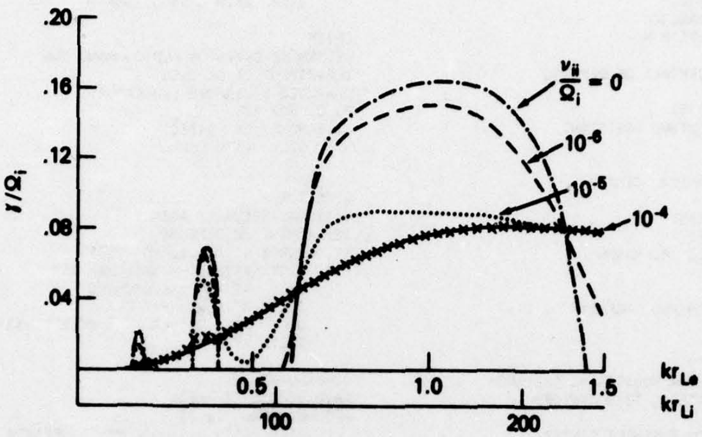

Fig. 1 — Plot of γ/Ω_i vs. kr_{Le} and kr_{Li} for $m_i = 16\ m_p$, $T_e = T_i$, $\omega_{pe}/\Omega_e = 10.0$, $V_{di}/V_i = 0.037$ and $\nu_{ii}/\Omega_i = 0.0$ (— · — · —), 10^{-6} (— — — —), 10^{-5} (· · · · · ·) and 10^{-4} (× × × × ×). The solid line represents the lower-hybrid-drift instability.

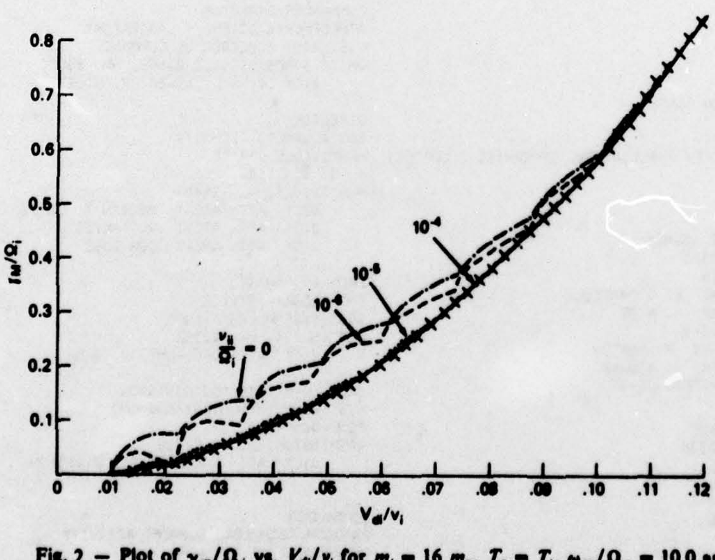

Fig. 2 — Plot of γ_m/Ω_i vs. V_{di}/v_i for $m_i = 16\ m_p$, $T_e = T_i$, $\omega_{pe}/\Omega_e = 10.0$ and $\nu_{ii}/\Omega_i = 0.0$ (— · — · —), 10^{-6} (— — — —), 10^{-5} (· · · · · ·) and 10^{-4} (× × × × ×). The solid line represents the lower-hybrid-drift instability.

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400 ARMY/NAVY DRIVE
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UNIVERSITY OF PITTSBURGH
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